

# Inflation, the $\mu$ Problem and Maximal $\nu_\mu - \nu_\tau$ Mixing

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## Abstract

A supersymmetric model based on a left-right symmetric gauge group, which ‘naturally’ leads to hybrid inflation, is studied. It is shown that the  $\mu$  problem can be easily solved in this model. The observed baryon asymmetry of the universe is produced via a primordial leptogenesis. For masses of  $\nu_\mu, \nu_\tau$  from the small angle MSW resolution of the solar neutrino problem and SuperKamiokande, maximal  $\nu_\mu - \nu_\tau$  mixing can be achieved. The required values of the relevant parameters are, however, quite small.

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A suitable framework for hybrid inflation [1] is provided [2] by a moderate extension of the minimal supersymmetric standard model (MSSM) based on a left-right symmetric gauge group. The inflaton is associated with the breaking of  $SU(2)_R$  and consists of a gauge singlet and a pair of  $SU(2)_R$  doublets. The  $\mu$  problem of MSSM can be resolved [3], in this scheme, by introducing [2,3] a trilinear superpotential coupling of the gauge singlet inflaton to the electroweak higgs doublets. It has been shown [3] that, in the presence of gravity-mediated supersymmetry breaking, this gauge singlet acquires a vacuum expectation value (vev) and consequently generates, through its coupling to the electroweak higgs superfields, the  $\mu$  term of MSSM.

The inflaton system, after the end of inflation, predominantly decays into electroweak higgs superfields and ‘reheats’ the universe. Moreover, its subdominant decay into right handed neutrinos provides a mechanism [4] for baryogenesis via a primordial leptogenesis. We solve [5] the evolution equations of the inflaton system and estimate the ‘reheat’ temperature. The process of baryogenesis is considered and its consequences on  $\nu_\mu - \nu_\tau$  mixing are analyzed [5]. For masses of  $\nu_\mu, \nu_\tau$  which are consistent with the small angle MSW resolution of the solar neutrino problem and the recent results of the SuperKamiokande experiment [6], we examine whether maximal  $\nu_\mu - \nu_\tau$  mixing can be achieved.

Let us consider a supersymmetric model based on the left-right symmetric gauge group  $G_{LR} = SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ . The breaking of  $SU(2)_R \times U(1)_{B-L}$  to  $U(1)_Y$  is achieved by the renormalizable superpotential

$$W = \kappa S(l^c \bar{l}^c - M^2) , \quad (1)$$

where  $S$  is a gauge singlet chiral superfield and  $l^c, \bar{l}^c$  is a conjugate pair of  $SU(2)_R$  doublet chiral superfields which acquire superheavy vacuum expectation values (vevs) of magnitude  $M$ . The parameters  $\kappa$  and  $M$  can be made positive by phase redefinitions.

It is well-known [2,7,8] that  $W$  in Eq.(1) leads ‘naturally’ to hybrid inflation [1]. This means that a) there is no need for ‘tiny’ coupling constants, b) the superpotential is the most general one allowed by the symmetry, c) supersymmetry guarantees that radiative corrections do not invalidate inflation, but rather provide a slope along the inflationary trajectory which drives the inflaton towards the supersymmetric vacua, and d) supergravity corrections can be brought under control so as to leave inflation intact.

The  $\mu$  problem can be resolved [3] by introducing the extra superpotential coupling

$$\delta W = \lambda S h^2 = \lambda S \epsilon^{ij} h_i^{(1)} h_j^{(2)} , \quad (2)$$

where the chiral electroweak higgs superfield  $h = (h^{(1)}, h^{(2)})$  belongs to a bidoublet  $(2, 2)_0$  representation of  $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$  and the parameter  $\lambda$  can again be made positive by suitable phase redefinitions. After gravity-mediated supersymmetry breaking,  $S$  acquires a vev which generates a  $\mu$  term [3].

The scalar potential which results from the superpotential terms in Eqs.(1) and (2) is (for simplicity, we take canonical Kähler potential):

$$V = |\kappa l^c \bar{l}^c + \lambda h^2 - \kappa M^2|^2 + (m_{3/2}^2 + \kappa^2 |\bar{l}^c|^2 + \kappa^2 |l^c|^2 + \lambda^2 |h|^2) |S|^2 + m_{3/2}^2 (|\bar{l}^c|^2 + |l^c|^2 + |h|^2) + \left( A m_{3/2} S (\kappa l^c \bar{l}^c + \lambda h^2 - \kappa M^2) + 2 \kappa m_{3/2} M^2 S + \text{h.c.} \right) , \quad (3)$$

where  $m_{3/2}$  is the universal scalar mass (gravitino mass) and  $A$  the universal coefficient of the trilinear soft terms. For exact supersymmetry ( $m_{3/2} \rightarrow 0$ ), the vacua are [3] at

$$S = 0, \quad \kappa l^c \bar{l}^c + \lambda h^2 = \kappa M^2, \quad l^c = e^{i\phi} \bar{l}^{c*} \quad h_i^{(1)} = e^{i\theta} \epsilon_{ij} h^{(2)j*}, \quad (4)$$

where the last two conditions arise from the requirement of D flatness. We see that there is a twofold degeneracy of the vacuum which is lifted by supersymmetry breaking. We get two degenerate (up to  $m_{3/2}^4$ ) ground states ( $\kappa \neq \lambda$ ): the desirable ('good') vacuum at  $h = 0$  and  $l^c \bar{l}^c = M^2$  and the undesirable ('bad') one at  $h \neq 0$  and  $l^c \bar{l}^c = 0$ . They are separated by a potential barrier of order  $M^2 m_{3/2}^2$ .

To leading order in supersymmetry breaking, the term of the potential  $V$  in Eq.(3) proportional to  $A$  vanishes, but a destabilizing tadpole term for  $S$  remains:

$$2 \kappa m_{3/2} M^2 S + \text{h.c.} . \quad (5)$$

This term together with the mass term of  $S$  (evaluated at the 'good' vacuum) give  $\langle S \rangle \approx -m_{3/2}/\kappa$  which, substituted in Eq.(2), generates [3] a  $\mu$  term with

$$\mu = \lambda \langle S \rangle \approx -\frac{\lambda}{\kappa} m_{3/2} . \quad (6)$$

Thus, coupling  $S$  to the higgses can lead to the resolution of the  $\mu$  problem.

The model can be extended [3] to include matter fields too. The superpotential has the most general form respecting the  $G_{LR}$  gauge symmetry and a global  $U(1)$  R-symmetry. Baryon number is automatically implied by this R-symmetry to all orders in the superpotential, thereby guaranteeing the stability of proton.

The model has [2,3] a built-in inflationary trajectory parametrized by  $|S|$ ,  $|S| > S_c = M$  for  $\lambda > \kappa$  (see below). All other fields vanish on this trajectory. The  $F_S$  term is constant providing a constant tree level vacuum energy density  $\kappa^2 M^4$ , which is responsible for inflation. One-loop radiative corrections (from the mass splitting in the supermultiplets  $l^c$ ,  $\bar{l}^c$  and  $h$ ) generate a logarithmic slope [8] along the inflationary trajectory which drives the inflaton toward the minimum. For  $|S| \leq S_c = M$ , the  $l^c$ ,  $\bar{l}^c$  components become tachyonic and the system evolves towards the ‘good’ supersymmetric minimum at  $h = 0$ ,  $l^c = \bar{l}^c = M$  (for  $\kappa > \lambda$ ,  $h$  is destabilized first and the system would have evolved towards the ‘bad’ minimum at  $h \neq 0$ ,  $l^c = \bar{l}^c = 0$ ). For all values of the parameters considered here, inflation continues at least till  $|S|$  approaches the instability at  $|S| = S_c$  as one deduces from the slow roll conditions [8,9]. The cosmic microwave quadrupole anisotropy can be calculated [8] by standard methods and turns out to be

$$\left(\frac{\delta T}{T}\right)_Q \approx \frac{32\pi^{5/2}}{3\sqrt{5}} \left(\frac{M}{M_P}\right)^3 \kappa^{-1} x_Q^{-1} \Lambda(x_Q)^{-1}, \quad (7)$$

where  $M_P = 1.22 \times 10^{19}$  GeV is the Planck scale and

$$\begin{aligned} \Lambda(x) = & \left(\frac{\lambda}{\kappa}\right)^3 \left[ \left(\frac{\lambda}{\kappa}x^2 - 1\right) \ln\left(1 - \frac{\kappa}{\lambda}x^{-2}\right) + \left(\frac{\lambda}{\kappa}x^2 + 1\right) \ln\left(1 + \frac{\kappa}{\lambda}x^{-2}\right) \right] \\ & + (x^2 - 1) \ln(1 - x^{-2}) + (x^2 + 1) \ln(1 + x^{-2}), \end{aligned} \quad (8)$$

with  $x = |S|/S_c$  and  $S_Q$  being the value of  $|S|$  when the present horizon scale crossed outside the inflationary horizon. The number of e-foldings experienced by the universe between the time the quadrupole scale exited the horizon and the end of inflation is

$$N_Q \approx 32\pi^3 \left(\frac{M}{M_P}\right)^2 \kappa^{-2} \int_1^{x_Q^2} \frac{dx^2}{x^2} \Lambda(x)^{-1}. \quad (9)$$

The spectral index of density perturbations turns out to be very close to unity.

After reaching the instability at  $|S| = S_c$ , the system continues [10] inflating for another e-folding or so reducing its energy density by a factor of about 2 – 3. It then

rapidly settles into a regular oscillatory phase about the vacuum. Parametric resonance is safely ignored in this case [10]. The inflaton (oscillating system) consists of the two complex scalar fields  $S$  and  $\theta = (\delta\phi + \delta\bar{\phi})/\sqrt{2}$ , where  $\delta\phi = \phi - M$ ,  $\delta\bar{\phi} = \bar{\phi} - M$ , with mass  $m_{infl} = \sqrt{2}\kappa M$  ( $\phi, \bar{\phi}$  are the neutral components of  $l^c, \bar{l}^c$ ).

The scalar fields  $S$  and  $\theta$  predominantly decay into electroweak higgsinos and higgses respectively with a common decay width  $\Gamma_h = (1/16\pi)\lambda^2 m_{infl}$ , as one can easily deduce from the couplings in Eqs.(1) and (2). Note, however, that  $\theta$  can also decay to right handed neutrinos  $\nu^c$  through the nonrenormalizable superpotential term

$$\frac{M_{\nu^c}}{2M^2} \bar{\phi} \bar{\phi} \nu^c \nu^c , \quad (10)$$

allowed by the gauge and R- symmetries of the model [2,3]. Here,  $M_{\nu^c}$  denotes the Majorana mass of the relevant  $\nu^c$ . The scalar  $\theta$  decays preferably into the heaviest  $\nu^c$  with  $M_{\nu^c} \leq m_{infl}/2$ . The decay rate is given by

$$\Gamma_{\nu^c} \approx \frac{1}{16\pi} \kappa^2 m_{infl} \alpha^2 (1 - \alpha^2)^{1/2} , \quad (11)$$

where  $0 \leq \alpha = 2M_{\nu^c}/m_{infl} \leq 1$ . The subsequent decay of these  $\nu^c$  's produces a primordial lepton number [4] which is then partially converted to the observed baryon asymmetry of the universe through electroweak sphaleron effects.

The energy densities  $\rho_S, \rho_\theta$ , and  $\rho_r$  of the oscillating fields  $S, \theta$ , and the 'new' radiation produced by their decay to higgsinos, higgses and  $\nu^c$  's are controlled by the equations:

$$\dot{\rho}_S = -(3H + \Gamma_h)\rho_S , \quad \rho_\theta(t) = \rho_S(t)e^{-\Gamma_{\nu^c}(t-t_0)} , \quad (12)$$

$$\dot{\rho}_r = -4H\rho_r + \Gamma_h\rho_S + (\Gamma_h + \Gamma_{\nu^c})\rho_\theta , \quad (13)$$

where

$$H = \frac{\sqrt{8\pi}}{\sqrt{3}M_P} (\rho_S + \rho_\theta + \rho_r)^{1/2} \quad (14)$$

is the Hubble parameter and overdots denote derivatives with respect to cosmic time  $t$ . The cosmic time at the onset of oscillations is taken  $t_0 \approx 0$ . The initial values of the various energy densities are taken to be  $\rho_S(t_0) = \rho_\theta(t_0) \approx \kappa^2 M^4/6$ ,  $\rho_r(t_0) = 0$ . The 'reheat' temperature  $T_r$  is calculated from the equation

$$\rho_S + \rho_\theta = \rho_r = \frac{\pi^2}{30} g_* T_r^4, \quad (15)$$

where the effective number of massless degrees of freedom is  $g_*=228.75$  for MSSM.

The lepton number density  $n_L$  produced by the  $\nu^c$  's satisfies the evolution equation:

$$\dot{n}_L = -3Hn_L + 2\epsilon\Gamma_{\nu^c}n_\theta, \quad (16)$$

where  $\epsilon$  is the lepton number produced per decaying right handed neutrino and the factor of 2 in the second term of the rhs comes from the fact that we get two  $\nu^c$  's for each decaying scalar  $\theta$  particle. The 'asymptotic' ( $t \rightarrow 0$ ) lepton asymmetry turns out to be

$$\frac{n_L(t)}{s(t)} \sim 3 \left( \frac{15}{8} \right)^{1/4} \pi^{-1/2} g_*^{-1/4} m_{infl}^{-1} \frac{\epsilon \Gamma_{\nu^c}}{\Gamma_h + \Gamma_{\nu^c}} \rho_r^{-3/4} \rho_S e^{\Gamma_h t}. \quad (17)$$

For MSSM spectrum between 100 GeV and  $M$ , the observed baryon asymmetry is then given [11] by  $n_B/s = -(28/79)(n_L/s)$ . It is, however, important to ensure that the primordial lepton asymmetry is not erased by lepton number violating  $2 \rightarrow 2$  scattering processes at all temperatures between  $T_r$  and 100 GeV. This requirement gives [11]  $m_{\nu_\tau} \lesssim 10$  eV which is readily satisfied in our case (see below).

Assuming hierarchical light neutrino masses, we take  $m_{\nu_\mu} \approx 2.6 \times 10^{-3}$  eV which is the central value of the  $\mu$ -neutrino mass coming from the small angle MSW resolution of the solar neutrino problem [12]. The  $\tau$ -neutrino mass will be restricted by the atmospheric anomaly [6] in the range  $3 \times 10^{-2}$  eV  $\lesssim m_{\nu_\tau} \lesssim 11 \times 10^{-2}$  eV. Recent analysis [13] of the results of the CHOOZ experiment shows that the oscillations of solar and atmospheric neutrinos decouple. We thus concentrate on the two heaviest families ignoring the first one. Under these circumstances, the lepton number generated per decaying  $\nu^c$  is [9,14]

$$\epsilon = \frac{1}{8\pi} g \left( \frac{M_3}{M_2} \right) \frac{c^2 s^2 \sin 2\delta (m_3^{D\ 2} - m_2^{D\ 2})^2}{|\langle h^{(1)} \rangle|^2 (m_3^{D\ 2} s^2 + m_2^{D\ 2} c^2)}, \quad (18)$$

where  $g(r) = r \ln(1 + r^{-2})$ ,  $|\langle h^{(1)} \rangle| \approx 174$  GeV,  $c = \cos \theta$ ,  $s = \sin \theta$ , and  $\theta$  ( $0 \leq \theta \leq \pi/2$ ) and  $\delta$  ( $-\pi/2 \leq \delta < \pi/2$ ) are the rotation angle and phase which diagonalize the Majorana mass matrix of  $\nu^c$  's with eigenvalues  $M_2, M_3$  ( $\geq 0$ ). The 'Dirac' mass matrix of the neutrinos is considered diagonal with eigenvalues  $m_2^D, m_3^D$  ( $\geq 0$ ).

For the range of parameters considered here, the scalar  $\theta$  decays into the second heaviest right handed neutrino with mass  $M_2$  ( $< M_3$ ) and, thus,  $M_{\nu^c}$  in Eqs.(10) and

(11) should be identified with  $M_2$ . Moreover,  $M_3$  turns out to be bigger than  $m_{infl}/2$  as it should. We will denote the two positive eigenvalues of the light neutrino mass matrix by  $m_2$  ( $=m_{\nu_\mu}$ ),  $m_3$  ( $=m_{\nu_\tau}$ ) with  $m_2 \leq m_3$ . All the quantities here (masses, rotation angles and phases) are ‘asymptotic’ (defined at the grand unification scale  $M_{GUT}$ ).

The determinant and the trace invariance of the light neutrino mass matrix imply [14] two constraints on the (asymptotic) parameters which take the form:

$$m_2 m_3 = \frac{(m_2^D m_3^D)^2}{M_2 M_3}, \quad (19)$$

$$m_2^2 + m_3^2 = \frac{(m_2^{D^2} c^2 + m_3^{D^2} s^2)^2}{M_2^2} + \frac{(m_3^{D^2} c^2 + m_2^{D^2} s^2)^2}{M_3^2} + \frac{2(m_3^{D^2} - m_2^{D^2})^2 c^2 s^2 \cos 2\delta}{M_2 M_3}. \quad (20)$$

The  $\mu - \tau$  mixing angle  $\theta_{23}$  ( $=\theta_{\mu\tau}$ ) lies [14] in the range

$$|\varphi - \theta^D| \leq \theta_{23} \leq \varphi + \theta^D, \text{ for } \varphi + \theta^D \leq \pi/2,$$

$$|\varphi - \theta^D| \leq \theta_{23} \leq \pi - \varphi - \theta^D, \text{ for } \varphi + \theta^D \geq \pi/2, \quad (21)$$

where  $\varphi$  ( $0 \leq \varphi \leq \pi/2$ ) is the rotation angle which diagonalizes the light neutrino mass matrix, and  $\theta^D$  ( $0 \leq \theta^D \leq \pi/2$ ) is the ‘Dirac’ (unphysical) mixing angle in the  $2 - 3$  leptonic sector defined in the absence of the Majorana masses of the  $\nu^c$ ’s.

Assuming approximate  $SU(4)_c$  symmetry, we get the asymptotic (at  $M_{GUT}$ ) relations:

$$m_2^D \approx m_c, \quad m_3^D \approx m_t, \quad \sin \theta^D \approx |V_{cb}|. \quad (22)$$

Renormalization effects, for MSSM spectrum and  $\tan \beta \approx m_t/m_b$ , are incorporated [14] by substituting in the above formulas the values:  $m_2^D \approx 0.23$  GeV,  $m_3^D \approx 116$  GeV and  $\sin \theta^D \approx 0.03$ . Also,  $\tan^2 2\theta_{23}$  increases by about 40% from  $M_{GUT}$  to  $M_Z$ .

We take a specific MSSM framework [15] where the three Yukawa couplings of the third generation unify ‘asymptotically’ and, thus,  $\tan \beta \approx m_t/m_b$ . We choose the universal scalar mass (gravitino mass)  $m_{3/2} \approx 290$  GeV and the universal gaugino mass

$M_{1/2} \approx 470$  GeV. These values correspond [16] to  $m_t(m_t) \approx 166$  GeV and  $m_A$  (the tree level CP-odd scalar higgs mass)  $= M_Z$ . The ratio  $\lambda/\kappa$  is evaluated [17] from

$$\frac{\lambda}{\kappa} = \frac{|\mu|}{m_{3/2}} \approx \frac{M_{1/2}}{m_{3/2}} \left(1 - \frac{Y_t}{Y_f}\right)^{-3/7} \approx 3.95, \quad (23)$$

where  $Y_t = h_t^2 \approx 0.91$  is the square of the top-quark Yukawa coupling and  $Y_f \approx 1.04$  is the weak scale value of  $Y_t$  corresponding to ‘infinite’ value at  $M_{GUT}$ .

Eqs.(7)-(9) can now be solved, for  $(\delta T/T)_Q \approx 6.6 \times 10^{-6}$  from COBE,  $N_Q \approx 50$  and any value of  $x_Q > 1$ . Eliminating  $x_Q$ , we obtain  $M$  as a function of  $\kappa$  depicted in Fig.1. The evolution Eqs.(12)-(14) are solved for each value of  $\kappa$ . The parameter  $\alpha^2$  in Eq.(11) is taken equal to  $2/3$ . This choice maximizes the decay width of the inflaton to  $\nu^c$ ’s and, thus, the subsequently produced lepton asymmetry. The ‘reheat’ temperature is then calculated from Eq.(15) for each value of  $\kappa$ . The result is again depicted in Fig.1.

We next evaluate the lepton asymmetry. We first take  $m_{\nu_\tau} \approx 7 \times 10^{-2}$  eV, the central value from SuperKamiokande [6]. The mass of the second heaviest  $\nu^c$ , into which the scalar  $\theta$  decays partially, is given by  $M_2 = M_{\nu^c} = \alpha m_{infl}/2$  and  $M_3$  is found from the ‘determinant’ condition in Eq.(19). The ‘trace’ condition in Eq.(20) is then solved for  $\delta(\theta)$  which is subsequently substituted in Eq.(18) for  $\epsilon$ . The leptonic asymmetry as a function of the angle  $\theta$  can be found from Eq.(17). For each value of  $\kappa$ , there are two values of  $\theta$  satisfying the low deuterium abundance constraint  $\Omega_B h^2 \approx 0.025$ . (These values of  $\theta$  turn out to be quite insensitive to the exact value of  $n_B/s$ .) The corresponding  $\varphi$ ’s are then found and the allowed region of the mixing angle  $\theta_{\mu\tau}$  in Eq.(21) is determined for each  $\kappa$ . Taking into account renormalization effects and superimposing all the permitted regions, we obtain the allowed range of  $\sin^2 2\theta_{\mu\tau}$  as a function of  $\kappa$ , shown in Fig.2. We observe that maximal mixing ( $\sin^2 2\theta_{\mu\tau} \approx 1$ ) is achieved for  $1.5 \times 10^{-6} \lesssim \kappa \lesssim 1.8 \times 10^{-6}$ . Also,  $\sin^2 2\theta_{\mu\tau} \gtrsim 0.8$  [6] corresponds to  $1.2 \times 10^{-6} \lesssim \kappa \lesssim 3.4 \times 10^{-6}$ .

We repeated the above analysis for all values of  $m_{\nu_\tau}$  allowed by SuperKamiokande. The allowed regions in the  $m_{\nu_\tau} - \kappa$  plane for maximal  $\nu_\mu - \nu_\tau$  mixing (bounded by the solid lines) and  $\sin^2 2\theta_{\mu\tau} \gtrsim 0.8$  (bounded by the dotted lines) are shown in Fig.3. Notice that, for  $\sin^2 2\theta_{\mu\tau} \gtrsim 0.8$ ,  $\kappa \approx (0.9 - 7.5) \times 10^{-6}$  which is rather small. (Fortunately, supersymmetry protects it from radiative corrections.) The corresponding values of  $M$  and  $T_r$  can be read from Fig.1. We find  $1.3 \times 10^{15}$  GeV  $\lesssim M \lesssim 2.7 \times 10^{15}$  GeV and



$10^7 \text{ GeV} \lesssim T_r \lesssim 3.2 \times 10^8 \text{ GeV}$ . We observe that  $M$  turns out to be somewhat smaller than the MSSM unification scale  $M_{GUT}$ . (It is anticipated that  $G_{LR}$  is embedded in a grand unified theory.) The reheat temperature, however, satisfies the gravitino constraint ( $T_r \lesssim 10^9 \text{ GeV}$ ). Note that, for the values of the parameters chosen here, the lightest supersymmetric particle (LSP) is [16] an almost pure bino with mass  $m_{LSP} \approx 0.43 M_{1/2} \approx 200 \text{ GeV}$  [18] and can, in principle, provide the cold dark matter of the universe. On the contrary, there is no hot dark matter candidate, in the simplest scheme.

In conclusion, we have shown that, in a supersymmetric model based on a left-right symmetric gauge group and leading ‘naturally’ to hybrid inflation, the  $\mu$  problem can be easily solved. The observed baryon asymmetry of the universe is produced via a primordial leptogenesis. For masses of  $\nu_\mu, \nu_\tau$  from the small angle MSW resolution of the solar neutrino puzzle and SuperKamiokande, maximal  $\nu_\mu - \nu_\tau$  mixing can be achieved. The required values of the coupling constant  $\kappa$  are, however, quite small ( $\sim 10^{-6}$ ).

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# FIGURES

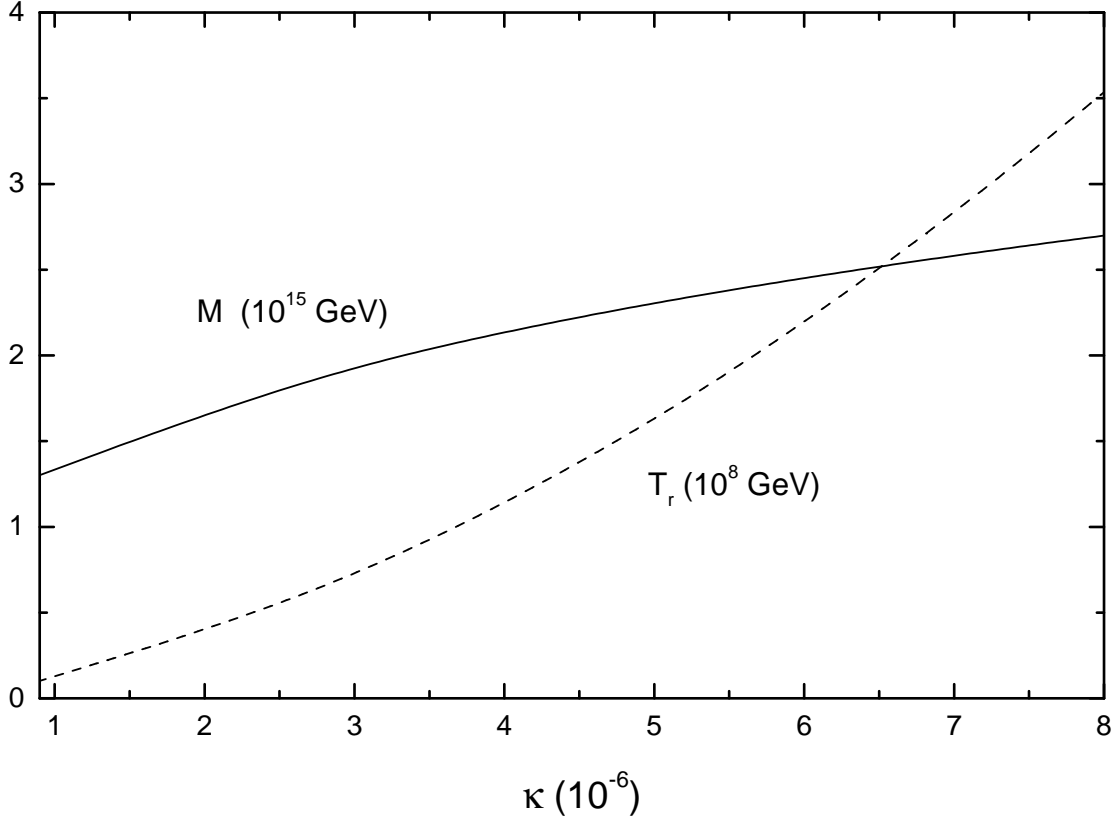


FIG. 1. The mass scale  $M$  (solid line) and the reheating temperature  $T_r$  (dashed line) as functions of  $\kappa$ .

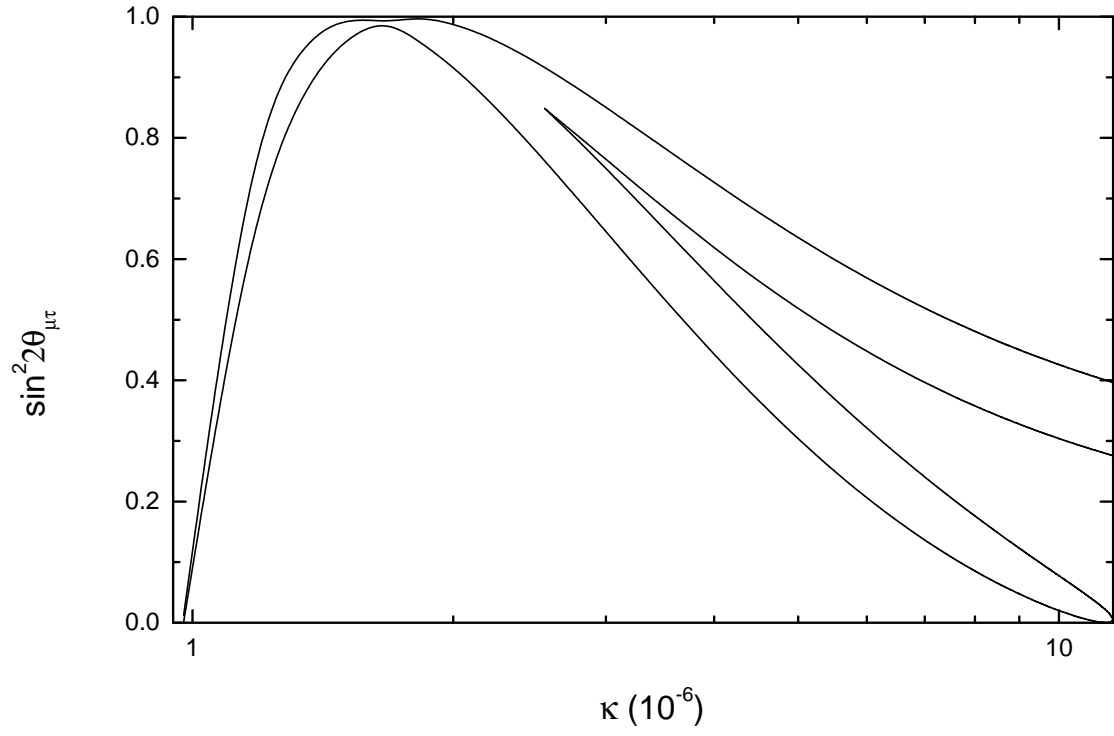


FIG. 2. The allowed region (bounded by the solid lines) in the  $\kappa - \sin^2 2\theta_{\mu\tau}$  plane for  $m_{\nu_\mu} \approx 2.6 \times 10^{-3}$  eV and  $m_{\nu_\tau} \approx 7 \times 10^{-2}$  eV.

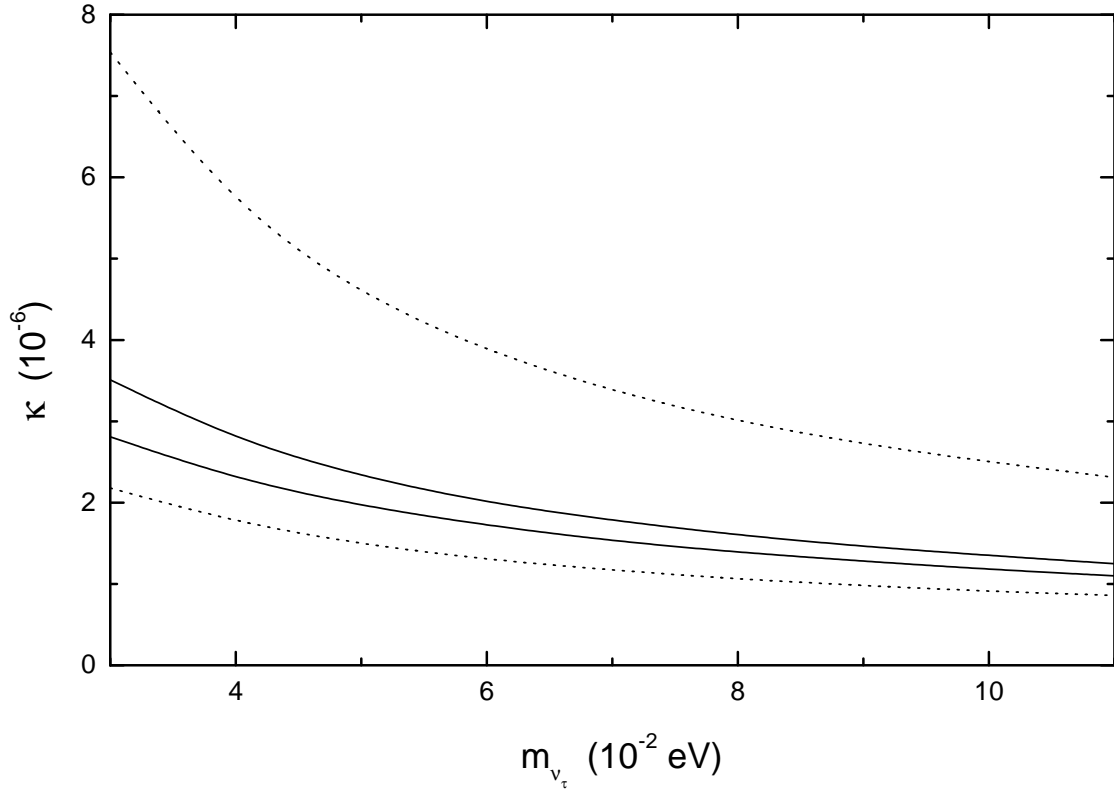


FIG. 3. The regions on the  $m_{\nu_\tau} - \kappa$  plane corresponding to maximal  $\nu_\mu - \nu_\tau$  mixing (bounded by the solid lines) and  $\sin^2 2\theta_{\mu\tau} \gtrsim 0.8$  (bounded by the dotted lines). Here we consider the range  $3 \times 10^{-2} \text{ eV} \lesssim m_{\nu_\tau} \lesssim 11 \times 10^{-2} \text{ eV}$  ( $m_{\nu_\mu} \approx 2.6 \times 10^{-3} \text{ eV}$ ).